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Transformations of Gaussian Light Beams Caused by Reflection in FEL Resonators

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The transformation of Gaussian radiation beams caused by reflection off mirrors is an important issue for FELs operating as oscillators. The reflected radiation from a single incident Gaussian mode will contain other modes due to the finite mirror size, the deflection of the beam and mismatches in the curvature. A method for analytic computation of the reflection matrix is developed by taking the convolution of the source function at the surface of the mirror with the paraxial propagator. The mirror surface that reflects spherical incoming wavefronts into spherical outgoing is determined analytically. Integral expressions for the reflection coefficients R_{pq}^{mn} for any incoming mode u_{mn} into the outgoing u_{pq} are obtained as functions of the deflection angle ϕ, the reflected beam waist w_0 and the mirror size. Analytic expansions of the reflection matrix in powers of the diffraction angle $\theta_d << 1$ are then obtained in both Hermite and Laguerre representations. The

phi   sub ϕ   u sub pq

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theta sub d

R sup mn sub pq   u sub mn

19. ABSTRACTS (Continued)

cross-coupling among vector components of the radiation field, caused by the curvature of the mirror surface, is included. It is shown that the lowest order contribution to the off-diagonal matrix elements is caused by the finite mirror size. The effects of the mirror curvature and the deflection of the light beam enter the reflection matrix as first order corrections in ϑ_{a}. theta sub a, [illegible]

Free Electron Lasers,
Oscillators
Optical Resonators
[illegible] Matrix (U)

CONTENTS

TRANSFORMATIONS OF GAUSSIAN LIGHT BEAMS CAUSED BY REFLECTION IN FEL RESONATORS

I. INTRODUCTION

Free Electron Lasers (FELs) operating as oscillators[1-7] require the trapping of light pulses between systems of mirrors (resonators).[8,9] These pulses are repeatedly amplified via synchronous interaction with electron pulses passing through the wiggler. The radiation produced by the stimulated emission is confined within a narrow cone along the beam axis. Therefore, the vector potential can be represented as a superposition of Gaussian modes. Those are the free space eigenmodes $\mathbf{A}_{mn}(r) = \mathbf{e}_{mn}A_{mn}(r)e^{ikz}$ where $\mathbf{e}_{mn}$ is the polarization vector, of the paraxial equation,[10]

$$\nabla_{\perp}^{2}A - 2ik \frac{\partial A}{\partial z} = 0. \tag{1}$$

Equation (1) is the $k_{\perp} << k = \omega/c$ limit of the exact wave equation. The simplest oscillator configuration is that of an open resonator with two opposed identical mirrors. The vacuum eigenmodes for this arrangement are also expressed in terms of the paraxial eigenmodes. Their detailed structure can be described in terms of either Gaussian-Hermite functions in rectangular coordinates, or Laguerre functions in polar coordinates. In both representations all the eigenmodes with given wave number k are characterized by two independent parameters: the waist $w = (2b/k)^{1/2}$ and the curvature of the wave front $1/R = z/(z^2 + b^2)$, where z is the distance from the waist position and b is the Rayleigh length (Fig. 1).

The electron beam is an optically active medium that alters the characteristic parameters of the radiation after each passage. During the build-up period the modal content and the structure of the light pulses inside the oscillator will change. A numerical method has been developed recently optimizing the representation for the amplified radiation. In the source dependent expansion[11,12] the waist size and the curvature of the elected modal basis is tailored according to the driving source term. That

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minimizes the number of modes required to describe the light beam. In general, the curvature and waist size of these modes does not match the curvature and waist of the vacuum eigenmodes for the resonator. Therefore, the transfer matrix for a given mirror must be known for arbitrary incoming modes. This need stems from computational as well as physical reasons. The knowledge of the cavity reflection matrix **R**, together with the gain matrix **G** through the wiggler, is necessary in determining the potential for steady state operation.

The study of the reflection matrix must include the effects of deflecting the light beam, in addition to finite mirror size and curvature mismatches. During high power operation, grazing mirror incidence may be necessary to avoid exceeding the dielectric breakdown limit for the reflecting surface. Also, in case of a high per-pass gain with optical guiding, the spot size for the reflected radiation could be much larger than the incoming. In two mirror resonators, the reflected radiation could then damage the wiggler. Therefore, ring resonators, including three or more mirrors, must be employed for the deflection and recirculation of the light pulses.

We are interested in cases when the reflected radiation remains focused along some direction $\hat{z}_o$ making an angle ϕ with the incoming $\hat{z}_i$. Then the reflected vector potential will also be expandable in free space eigenmodes $A_{pq}(r_o)$ of the paraxial equation in the new direction. The mirror surface generating focused reflection in the desired direction can not be arbitrary but must be appropriately defined. The angle of deflection ϕ will enter the equation defining the mirror surface. The other surface parameter, namely the curvature $1/R_m$, is a free parameter. It determines the curvature $1/R_o$ for the outgoing modes given the curvature $1/R_i$ of the incoming modes. In case of reflection by an arbitrary surface,

the scattered radiation cannot, in general, be covered by the paraxial modes that do not form a complete set in three dimensions.

A single incident mode $A_{mn}(r_i)$ will, in general, be partially reflected into different modes $A_{pq}(r_o)$ where $(m,n)\neq(p,q)$. This is caused by the deflection of the light beam, the finite size of the mirror and the curvature mismatches. Reflection into other modes will affect the interaction between the electron beam and the radiation in a number of ways. First, the rms radius of the light beam will change, affecting the matching beam condition. Second, the light pulse will spread axially because of dispersion among different modes, since the phase velocity depends on the modal number (m,n). Third, different phase shifts among the various modes during reflection may render these modes out of phase after a number of bouncings off the resonator. For the above reasons the fraction of radiation scattered into other modes will contribute to the losses in FEL oscillators.

The method for obtaining the reflection matrix is outlined in Sec. II. The definition of the appropriate mirror surface is given in Sec. III. In Sec. IV the integral expressions for the matrix elements are derived. An analytic expansion in powers of a small parameter (of the order of the diffraction angle) is given in the same section. Some limiting cases are examined in Sec. V. In Sec. VI the reflection of the fundamantal mode (0,0) is studied in detail. Section VII deals with cross-coupling effects among the vector components of the radiation.

II. OUTLINE OF THE METHOD

The free space eigenmodes $A_{mn}(r)$ of the paraxial wave equation have the general form

$$A_{mn}(r) = \frac{u_{mn}(r;W)}{\left(1 + \frac{z^2}{b^2}\right)^{1/2}} e^{i\left[kz + \frac{k(x^2 + y^2)}{2R(z)}\right]} e^{i\delta_{mn}(z)} . \tag{2}$$

The first exponential in (2) contains the rapidly varying phase on the wavelength scale $\lambda = 2\pi/k$. The wavefronts are spherical with radius of curvature $R(z)$ given by $1/R(z) = z/(z^2 + b^2)$. The spot size of the radiation envelope is $W(z) = w(1 + z^2/b^2)^{1/2}$, where $w = (2b/k)^{1/2}$ is the waist, and the distance z is measured from the position of the waist. The amplitude squared of the mode drops by 1/2 over a distance equal to the Rayleigh length b (also known as confocal parameter). Most of the radiation is confined within a cone parametrized by the diffraction angle $\theta_d = W/z \simeq (\lambda/b\pi)^{1/2}$. The structure of the amplitude profile $u_{mn}(r;W)$ depends on the elected coordinate system. $u_{mn}(r;W)$ contains the slow spatial variation equivalent to a small wave number perpendicular to the z-direction. Higher modes correspond to an increasing effective $k_\perp$, producing the slow phase shift expressed by the term $\exp\,[i\delta_{mn}(z)]$. For a given k, the mode is completely defined by the two independent parameters R and w (or any combination of two out of the four quantities R, w, z and b).

The geometry of the reflection is illustrated in Fig. 2. The subscripts i and o denote the coordinate system used for expressing incoming and outgoing modes. r_i is defined with the $\hat{z}_i$ axis along the direction of incidence and r_o has the $\hat{z}_o$ axis along the direction of reflection. The origins are displaced from the mirror center by l_i and l_o respectively, where l_i is the distance of the minimum waist w_i for the

incoming radiation and l_o is the distance of the minimum waist w_o for the reflected mode. A third coordinate system r_s with the origin at the mirror center and $\hat{z}_m$ aligned with the mirror axis will be useful in the computations. Underlined quantities $\underline{r}_i$, $\underline{r}_o$ and $\underline{r}_s$ stand for the mirror surface coordinates in each reference frame. The transformations among the various frames are defined by

$$\begin{aligned} x_i &= x_s \cos\frac{\phi}{2} - z_s \sin\frac{\phi}{2}, \\ y_i &= y_s, \\ z_i &= z_s \cos\frac{\phi}{2} + x_s \sin\frac{\phi}{2} + l_i, \end{aligned} \tag{3a}$$

$$\begin{aligned} x_o &= x_s \cos\frac{\phi}{2} + z_s \sin\frac{\phi}{2}, \\ y_o &= y_s, \\ z_o &= z_s \cos\frac{\phi}{2} - x_s \sin\frac{\phi}{2} + l_o. \end{aligned} \tag{3b}$$

We consider incoming radiation of given curvature and of arbitrary amplitude profile $A^i(r_i)$, consisting of various modes (m,n) with the same $R_i(z)$. If both incident and reflected radiation are expanded into eigenmodes,

$$A^i(\underline{r}_i) = e^{i\Phi_i(\underline{r}_i)} \sum_{m,n} c^i_{mn} \frac{u_{mn}(\underline{r}_i)}{\left[1 + \frac{z_i^2}{b_i^2}\right]^{1/2}} e^{i\delta_{mn}}, \tag{4a}$$

$$A^o(\underline{r}_o) = e^{i\Phi_o(\underline{r}_o)} \sum_{p,q} c^o_{pq} \frac{u_{pq}(\underline{r}_o)}{\left[1 + \frac{z_o^2}{b_o^2}\right]^{1/2}} e^{i\delta_{pq}}, \tag{4b}$$

where

$$\Phi_i(r) = k\left[z + \frac{x^2 + y^2}{2R_i(z)}\right], \tag{4c}$$

the relation among the incident and reflected expansion coefficients c^i_{mn}, c^o_{pq} is formulated by

$$\mathbf{c}^o = \mathbf{R}\,\mathbf{c}^i, \tag{5a}$$

or

$$c^o_{pq} = \sum_{m,n} R^{mn}_{pq}\, c^i_{mn}, \tag{5b}$$

where R^{mn}_{pq} are the elements of the reflection matrix $\mathbf{R}$.

We examine the case when the mirror dimensions ρ are much larger than the wavelength λ, $\lambda \ll \rho$ (otherwise diffraction rather than reflection would prevail). We also assume that the angle ζ subtended by the mirror $\zeta = \rho/R_m$, where R_m parametrizes the radius of curvature, is small, of the order of the diffraction angle θ_d, $\zeta \sim \theta_d \sim \varepsilon$. The ν-th component of the reflected vector potential at distance $|\mathbf{r}_o - \underline{\mathbf{r}}_o| \gg \lambda$ from the mirror surface S is then given by

$$A^o_{(\nu)}(\mathbf{r}_o) = -\frac{ik}{2\pi}\iint_S ds\, \frac{e^{ik|\mathbf{r}_o - \underline{\mathbf{r}}_o|}}{|\mathbf{r}_o - \underline{\mathbf{r}}_o|}\, A^s_{(\nu)}(\underline{\mathbf{r}}_o)\,(\hat{n}\cdot\hat{\delta r}). \tag{6}$$

In Eq. (6) $\hat{n}\cdot\hat{\delta r}$ is the obliqueness factor where $\hat{\delta r} = (\mathbf{r}_o-\underline{\mathbf{r}}_o)/|\mathbf{r}_o-\underline{\mathbf{r}}_o|$ and $\hat{n}$ is the normal unit vector to the reflecting surface. The surface element ds is given by $ds = \delta[z_o - f(x_o,y_o)]dx_o dy_o dz_o$ where $z_o = f(x_o,y_o)$ is the surface equation. Equation (6) is the convolution of a source term $A^s(\underline{\mathbf{r}}_o)$ at the mirror surface with the propagator $\exp(ik|\mathbf{r}_o-\underline{\mathbf{r}}_o|)/|\mathbf{r}_o-\underline{\mathbf{r}}_o|$, i.e., a superposition of spherical waves originating at S. The source term $\mathbf{A}^s(\underline{\mathbf{r}}_o)$ is specified from the incoming vector potential $\mathbf{A}^i(\underline{\mathbf{r}}_i)$ through the boundary conditions and the coordinate transformations (3). We will assume a perfectly conducting surface, where the incident and reflected fields are related by

$$\mathbf{A}^s = -\mathbf{A}^i + 2\,(\hat{n}\cdot\mathbf{A}^i)\,\hat{n}, \tag{7a}$$

and $\hat{n}$ is the normal unit vector to the reflecting surface. The second term in (7a) introduces a coupling among different vector components, caused by the mirror curvature. This cross-coupling is small and disappears in the plane mirror limit,

$$A^s_{(\nu)} = -\,A^i_{(\nu)}, \tag{7b}$$

where A^i and A^s are expressed in the incoming and outgoing coordinate systems respectively. Because of the linear superposition principle, Eq. (6), the cross-coupling contribution can be added separately, and will be deferred until Sec. VII. In the next three sections we will treat the reflected vector components as independent scalars, according to (7b), that corresponds to a phase shift by π during reflection. Most of the computations will be performed on the mirror surface. To simplify the notation from now on, we drop the bar (_) under the mirror coordinates $\underline{r}$. Subscripted quantities such as r_i, r_o, r_s will signify the mirror coordinates in each reference frame. Unsubscripted coordinates will denote the observation point in the reflected radiation frame of reference.

We seek cases when the reflected radiation propagates along z_o, contained within a cross section of dimensions $x,y \ll z-z_o$. The expansion $|r-r_o| \simeq (z-z_o)\,\{1 + [(x-x_o)^2 + (y-y_o)^2]/2(z-z_o)^2\}$ replaces the full propagator inside (6) with the paraxial propagator in that direction,

$$A^o(r) = \iint_S ds\; A(r_o)\;(\hat{n}\cdot\hat{\delta r})\;U_{-k}(r,r_o), \tag{8a}$$

where

$$U_{-k}(\mathbf{r}, \mathbf{r}_o) = \frac{ik}{2\pi} \frac{e^{-ik(z-z_o)}}{z-z_o} e^{-ik\frac{(x-x_o)^2 + (y-y_o)^2}{2(z-z_o)}} . \tag{8b}$$

Expression (8) is the approximation of the exact solution (6) to order $[(x-x_o)^2 + (y-y_o)^2]/2(z-z_o)^2 \sim \varepsilon^2$. It is valid provided the surface S produces focused reflection along the desired direction. Otherwise the paraxial limit will fail to encompass all the radiation contained in the original expression (6). The geometry of the mirror that is compatible with the above approximation will be obtained during the computation of the reflection matrix.

It is known that the profile of a given eigenmode $A_{mn}(x_o,y_o,z_o)$ at z_o is generated by the propagator $U_k(\mathbf{r},\mathbf{r}_o)$ acting on the mode $A_{mn}(x,y,0)$ at $z = 0$. The inverse propagator $U_{-k}(\mathbf{r},\mathbf{r}_o)$ therefore reproduces $A_{mn}(x,y,0)$ from $A_{mn}(x_o,y_o,z_o)$,

$$\iint_S dx_o dy_o \frac{u_{mn}(x_o,y_o,z_o)}{\left[1 + \frac{z_o^2}{b_o^2}\right]^{1/2}} e^{-ik\left[z_o + \frac{x_o^2 + y_o^2}{2R(z_o)}\right]} U_{-k}(r,r_o) = u_{mn}(x,y,0). \tag{9}$$

This suggests multiplying and dividing the integrand inside (8a) by $\exp[i\Phi(\mathbf{r}_o)] \,/\, [1 + z_o^2/b_o^2]^{1/2}$, recasting (8a) in the form,

$$A^o(\mathbf{r}) = \iint ds\; e^{i\Delta(\mathbf{r}_o)} S(\mathbf{r}_o)\, e^{-i\Phi_o(\mathbf{r}_o)}\; U_{-k}(\mathbf{r},\mathbf{r}_o), \tag{10}$$

where the source term $S(\mathbf{r}_o)$ is,

$$S(\mathbf{r}_o) = A^i[\mathbf{r}_i(\mathbf{r}_o)]\;(\hat{n} \cdot \hat{\delta r})\left[1 + \frac{z_o^2(\mathbf{r}_o)}{b_o^2}\right]^{1/2}, \tag{11}$$

and the phase term $\Delta(\mathbf{r}_o) = \Phi_i[\mathbf{r}_i(\mathbf{r}_o)] + \Phi_o(\mathbf{r}_o)$ is given by,

$$\Delta(\mathbf{r}_o) = k \left[z_i(\mathbf{r}_o) + z_o + \frac{x_i^2(\mathbf{r}_o) + y_i^2(\mathbf{r}_o)}{2R_i(\mathbf{r}_o)} + \frac{x_o^2 + y_o^2}{2R_o(\mathbf{r}_o)} \right]. \qquad (12)$$

The phase $\Delta(\mathbf{r}_o)$ depends on the angle ϕ through the coordinate transformations Eqs. (3).

The term $\exp[i\Delta(\mathbf{r}_o)]$ is varying rapidly, on the scale of the wavelength λ. Therefore, its convolution with the slowly varying source term over an arbitrary surface will be vanishingly small. In general, this corresponds to radiation scattering where only a small fraction of the incoming radiation is reflected along the considered direction ϕ. The integral (10) will be finite only when it is possible to satisfy the condition $\Delta(\mathbf{r}_o) \simeq$ constant over some surface S. If, in addition, S is much larger than λ, expression (10) will be finite only within a narrow angle $\delta\phi$ around ϕ. This guarantees that the reflected radiation remains focused along that direction. Therefore, a condition that the exact reflected radiation (6) be fully covered by the paraxial limit (10) is that

$$\Delta(\mathbf{r}_o) = \text{constant}, \qquad (13)$$

along the surface S. Accordingly, the optical path is the same along the rays connecting an incoming wave front with its mirror image (reflected) wave front.

Requirement (13) defines the appropriate mirror surface $z_o=f_o(x_o,y_o;\phi)$ for reflection in the elected direction. Assuming that f_o is found, we may express z_o in terms of x_o, y_o and use the constancy of $\Delta(\mathbf{r}_o)$ over S, reducing (10) into

$$A^o(\mathbf{r}) = \iint_S dx_o dy_o \; \sigma(x_o,y_o) \; e^{-i\Phi_o(x_o,y_o)} \; U_{-k}(\mathbf{r},\mathbf{r}_o). \qquad (14)$$

$\sigma(x_o,y_o) \equiv S[x_o,y_o,z_o(x_o,y_o)]$ is fully expanded in terms of $u_{mn}(x_o,y_o)$ that form a complete set in two dimensions,

$$\sigma(x_o,y_o) = \sum_{m,n} R^{mn}\, u_{mn}(x_o,y_o;W_o). \tag{15}$$

The expansion coefficients R^{mn} for Gaussian incoming radiation of arbitrary profile $\sigma(x_o,y_o)$ are given by

$$R^{mn} = \iint dx_o dy_o\, \sigma(x_o,y_o)\, u_{mn}(x_o,y_o;W_o) / \iint dx_o dy_o u_{mn}{}^2(x_o,y_o;W_o). \tag{16}$$

The radiation spot size W_o at the location of the mirror center is a free parameter, yet to be specified. Each choice of W_o generates an equivalent representation for $\sigma(x_o,y_o)$.

Upon substituting expansion (15) inside the integral (14) and using the property (9) for the inverse propagator U_{-k}, the reflected vector potential assumes the final form

$$A^o(x,y,0) = \sum_{m,n} R^{mn}\, u_{mn}(x,y;W_o), \tag{17}$$

where $W_o(z) = w_o\,(1 + z^2/b_o{}^2)^{1/2}$, $w_o = (2b_o/k)^{1/2}$. Expression (17) is a complete decomposition of the reflected radiation into paraxial eigenmodes for incident radiation of arbitrary profile. Therefore, condition (13) that defines the mirror surface is sufficient for the full reflection of paraxial (Gaussian) incoming light beams into paraxial beams only. The fraction of the electromagnetic flux incident on the mirror is conserved after reflection. If, on the other hand, (13) is seriously violated, the paraxial modes are inadequate to include all reflected radiation, and the incident flux is not conserved by expressions similar to (17).

III. MIRROR SURFACE.

To obtain the equation for S we express all quantities inside (12) in the mirror coordinate frame applying the transformations (3a) and (3b). Using the scaling $x_s/R_m \sim y_s/R_m \sim \varepsilon << 1$, $z_s/R_m \sim \varepsilon^2$ we obtain from (13)

$$z_s = - \frac{1}{2R_m \cos \frac{\phi}{2}} \left[x_s^2 \cos^2 \frac{\phi}{2} + y_s^2\right], \tag{18a}$$

where

$$\frac{1}{R_m} = \frac{1}{2R_o} + \frac{1}{2R_i}. \tag{18b}$$

Equation (18a) is the analytic expression for a paraboloid surface. R_m parametrizes the mirror curvature, being positive or negative for a convex or concave mirror respectively. The surface is reflection symmetric with $(zx)_s$ and $(zy)_s$ as the symmetry planes; there is no rotational symmetry around $\hat{z}_s$. Surface (18a) can also be approximated, to second order in $(x_s/R_m)^2$, $(y_s/R_m)^2$ by hyperboloids or ellipsoids defined by

$$\left(z_s - R_m \cos \frac{\phi}{2}\right)^2 - x_s^2 \cos^2 \frac{\phi}{2} - y_s^2 = R_m^2 \cos^2 \frac{\phi}{2}, \tag{19a}$$

$$\left(z_s + R_m \cos \frac{\phi}{2}\right)^2 + x_s^2 \cos^2 \frac{\phi}{2} + y_s^2 = R_m^2 \cos^2 \frac{\phi}{2}. \tag{19b}$$

All the surfaces become spherical in the limit of perpendicular incidence $\phi = 0$, and plane mirrors when $R_m \to \infty$. Using the definition of the curvature for the paraxial modes, Eq. (2), and the fact that R >> b in cases of interest, we obtain from (18b)

$$\frac{1}{R_o} = \frac{2}{R_m} - \frac{1}{R_i}. \tag{20}$$

Relation (20) defines the curvature of the reflected modes from the incoming mode curvature and the curvature of the mirror.

Equations (18)-(20) imply that

$$\Delta(\mathbf{r_s}) \equiv \Delta[\mathbf{r_i}(\mathbf{r_s}),\mathbf{r_o}(\mathbf{r_s})] = \text{const.} + O\left[k\rho \left(\frac{\rho}{R_m}\right)^2 \right], \tag{21}$$

where ρ parametrizes the mirror size. A more complicated surface equation (higher than quadratic in x, y, z) is required to improve the constancy to a higher order. In the next section the reflection matrix will be computed by expansion in powers of $W_o/R_m \simeq \rho/R_m$. Since $k\rho >> 1$, the approximation $\Delta(\mathbf{r_s})$ = constant is satisfactory for a first order expansion as long as $\rho/R_m \sim 1/k\rho$. In case that $\rho/R_m > 1/k\rho$, $\Delta(x_s,y_s)$ is a slowly varying function over S. Large mirrors require the inclusion of the phase slippage term $\exp[i\Delta(x_s,y_s)]$ next to the source term $\sigma(x_s,y_s)$ in Eq. (16).

The unit vector $\hat{n}$ normal to the mirror surface is given by

$$\hat{n} = \frac{\nabla f}{|\nabla f|} \simeq \cos\frac{\phi}{2}\,\frac{x_s}{R_m}\,\hat{x}_s + \frac{1}{\cos\frac{\phi}{2}}\,\frac{y_s}{R_m}\,\hat{y}_s + \left(1 + \frac{z_s}{R_m\cos\frac{\phi}{2}}\right)\hat{z}_s,$$

where f (x_s, y_s, z_s) is given by Eq. (18a).

IV. COMPUTATION OF THE REFLECTION MATRIX

According to the definition (5b), the R_{pq}^{mn} element of the reflection matrix **R** is obtained from the source term $\sigma_{pq}(x_o,y_o)$ inside (14) generated by a single incident eigenmode $A_{pq}[r_i(r_o)]$. The integration is performed in the mirror-aligned coordinates, taking advantage of the existing symmetries. The coordinates r_i and r_o, defining the incoming and outgoing wave functions, become explicit functions of x_s, y_s through the transformations (3). The surface equation (14a) is used to express z_s in terms of (x_s,y_s). The mirror boundary

$$x_s^{\,2}\cos^2\frac{\phi}{2} + y_s^{\,2} = \rho^2 \tag{22}$$

is defined by the intersection of the infinite surface (18a) with the plane $z_s = \text{const} = 2\rho^2\cos^2(\phi/2)/R_m$. After the above manipulations, the reflection matrix elements take the form

$$R_{pq}^{mn} = \iint_S dx_s dy_s \frac{\bar{u}_{mn}(x_s,y_s)\bar{u}_{pq}(x_s,y_s)}{\left[1 + \dfrac{l_o^{\,2}}{b_o^{\,2}}\right]^{1/2}} \left[\frac{1 + \dfrac{z_o^{\,2}(x_s,y_s)}{b_o^{\,2}}}{1 + \dfrac{z_i^{\,2}(x_s,y_s)}{b_i^{\,2}}}\right]^{1/2} e^{i\Delta(x_s,y_s)}$$

$$\times\, e^{i\delta_{pq}^{i}(x_s,y_s) - i\delta_{mn}^{o}(x_s,y_s)} \left[\cos\frac{\phi}{2}\left(1 - \frac{x_s}{R_m}\sin\frac{\phi}{2} - \frac{x_s^2\sin^2\frac{\phi}{2}}{R_m^{\,2}}\right)\right], \tag{23}$$

where

$$\bar{u}_{mn}(x_s,y_s) \equiv u_{mn}\left[x_o(x_s,y_s),\, y_s\right],\quad \bar{u}_{pq}(x_s,y_s) \equiv u_{pq}\left[x_i(x_s,y_s),\, y_s\right]. \tag{24}$$

Expression (23) is correct to order ρ^2/R_m^2.

It will be seen that **R**, as given by (23), depends on four parameters

$$\mathbf{R} \equiv \mathbf{R}(\phi,\ \alpha,\ \mu;\ \xi). \tag{25}$$

ϕ is the reflection angle shown in Fig. 2. α is the ratio of the incoming to the outgoing spot size at the mirror, $\alpha = W_i(l_i)/W_o(l_o)$. $\mu = \rho/W_o$ parametrizes the mirror size compared to the radiation spot size. $\xi = W_o/R_m$ scales as the diffraction angle $\theta_d \approx W_o/l_o$ multiplied by the curvature mismatch R_o/R_m between the mirror and the radiation wavefronts. The spot size W_o enters as a free parameter because only the curvature $1/R_o$ for the reflected modes is specified by the mirror geometry. Since many combinations of W_o and l_o apply to a given curvature according to paragraph Eq. (2), an additional selection rule for W_o is needed. Note that W_o does not have to match W_i. This is obvious in cases when the mirror size ρ is smaller that W_i. Each value of W_o defines a complete set of modes for the reflected radiation and an equivalent representation for **R**.

Parameters ϕ, α, and μ can be arbitrary. In most cases of interest, however, ξ is small, $\xi \ll 1$, of the same order as the diffraction angle θ_d. The analytic computation of the matrix elements is carried out by expanding the integral (23) in powers of ξ,

$$\mathbf{R} = \mathbf{R}(0) + \xi\ \mathbf{R}(1) + \xi^2\ \mathbf{R}(2). \tag{26}$$

Each representation of **R** is tied to the choice of the basis functions $u_{mn}(r)$. The eigenmodes $u_{mn}(r)$ are specified according to the coordinate geometry. In the next subsections we derive **R** in Gaussian-Hermite and Gaussian-Laguerre representations. For simplicity, it is assumed that $\Delta(x_s,y_s)$ in Eq. (23) is constant, i.e., $k\rho\ (\rho/R_m) \ll 1$.

(a). Gaussian-Hermite representation

In rectangular coordinates (x,y,z) the functions $u_{mn}(x,z;W)$ are given by

$$u_{mn}(x,y;W) = a_{mn}H_m\left(\frac{\sqrt{2}x}{W}\right) H_n\left(\frac{\sqrt{2}y}{W}\right) e^{-\frac{x^2+y^2}{W^2}} , \tag{27a}$$

where H_m, H_n are the Hermite polynomials and a_{mn} is a normalization factor, setting the total electromagnetic flux carried by the mode equal to unity,

$$a_{mn} = \frac{\sqrt{2}}{W}\left(\pi\ 2^{m+n}\ m!n!\right)^{-1/2} . \tag{27b}$$

The corresponding slow phase factor $\delta_{mn}(z)$ in Eq. (2) is

$$\delta_{mn}(z) = (m + n + 1)\tan^{-1}\left(\frac{z}{b}\right). \tag{27c}$$

Substituting inside (23), expanding in ξ and performing the integrations, Eqs. (23)-(26) yield

$$R^{mn}_{pq}(0) = C^{mn}_{pq}\ e^{i\psi^{mn}_{pq}}\ I^{mn}_{pq} , \tag{28a}$$

$$R^{mn}_{pq}(1) = C^{mn}_{pq}\ e^{i\psi^{mn}_{pq}}\ \tan\frac{\phi}{2}\left\{M^{mn}_{pq} + i\ N^{mn}_{pq}\right\} , \tag{28b}$$

where C^{mn}_{pq} is a normalization factor

$$C^{mn}_{pq} = \frac{W_o}{\pi W_i}\left(2^{m+n+p+q}\ m!n!p!q!\right)^{-1/2} , \tag{28c}$$

and the phase ψ^{mn}_{pq} is expressed by

$$\psi^{mn}_{pq} = (p+q+1)\tan^{-1}\left(\frac{l_i}{b_i}\right) - (m+n+1)\tan^{-1}\left(\frac{l_o}{b_o}\right) + k(l_i+l_o). \tag{28d}$$

$$I_{pq}^{mn} = \int_{-X_s}^{X_s} dX \int_{-Y_s}^{Y_s} dY\, H_p(\alpha X)\, H_q(Y)\, H_m(X)\, H_n(Y)\, e^{-\frac{\alpha^2+1}{2}(X^2+Y^2)} , \quad (29a)$$

$$M_{pq}^{mn} = \int_{-X_s}^{X_s} dX \int_{-Y_s}^{Y_s} dY\, H_p(\alpha X)\, H_q(Y)\, H_m(X)\, H_n(Y)\, e^{-\frac{\alpha^2+1}{2}(X^2+Y^2)}$$

$$\left\{-\frac{3}{2} X + \frac{1-\alpha^2}{\sqrt{2}} X(X^2+Y^2) - \left(\alpha \frac{H_p'(\alpha X)}{H_p(\alpha X)} - \frac{H_m'(X)}{H_m(X)}\right) \frac{\sqrt{2}}{4} (X^2+Y^2)\right\} , \quad (29b)$$

$$N_{pq}^{mn} = \int_{-X_s}^{X_s} dX \int_{-Y_s}^{Y_s} dY \frac{X}{\sqrt{2}}\, H_p(\alpha X)\, H_q(Y)\, H_m(X)\, H_n(Y)\, e^{-\frac{\alpha^2+1}{2}(X^2+Y^2)} . \quad (29c)$$

In the rescaled variables $X \simeq \cos\phi/2 \; \sqrt{2}x_s/W_o$, $Y = \sqrt{2}y_s/W_o$, the surface boundary is given by $X_s^{\,2} + Y_s^{\,2} = 2\rho^2/W_o^{\,2}$. The lowest terms can be computed directly. The matrix elements are computed, to first order in ξ, in Appendix A for large size mirror and $\alpha = 1$.

(b). Gaussian-Laguerre representation

In cylindrical coordinates (r,θ,z) where $\tan\theta = x/y$, $r = (x^2+y^2)^{1/2}$, $u_m^p(r,\theta;W)$ take the form

$$u_m^{\pm p}(r,\theta;W) = a_m^p \begin{pmatrix}\cos p\theta \\ \sin p\theta\end{pmatrix} \left(\frac{\sqrt{2}r}{W}\right)^p L_m^p\left(\frac{2r^2}{W^2}\right) e^{-\frac{1}{2}\frac{2r^2}{W^2}} , \quad (30a)$$

where $+p(-p)$ signifies cosine (sine) poloidal dependence, a_m^p is given by

$$a_m^p = \left(\frac{4}{\pi W^2}\right)^{1/2} \left(\frac{m!}{(m+p)!}\right)^{1/2} , \quad (30b)$$

and the L^p_m are the Laguerre polynomials. The corresponding slow phase $\delta^p_m(z)$ in Eq. (2) is,

$$\delta^p_m(z) = (2m + p + 1)\tan^{-1}\left(\frac{z}{b}\right). \tag{30c}$$

The transformations among polar coordinates representing the various reference frames are

$$r_i \approx r_s\left[1 - \sin^2\theta_s \sin^2\frac{\phi}{2} - 2\frac{z_s}{R_m}\sin\theta_s \sin\frac{\phi}{2}\cos\frac{\phi}{2}\right]^{1/2},$$

$$r_o = r_s\left[1 - \sin^2\theta_s \sin^2\frac{\phi}{2} + 2\frac{z_s}{R_m}\sin\theta_s \sin\frac{\phi}{2}\cos\frac{\phi}{2}\right]^{1/2}, \tag{31a}$$

$$\tan\theta_i = \cos\frac{\phi}{2}\tan\theta_s - \frac{z_s}{r_s}\frac{\sin\frac{\phi}{2}}{\cos\theta_s},$$

$$\tan\theta_o = \cos\frac{\phi}{2}\tan\theta_s + \frac{z_s}{r_s}\frac{\sin\frac{\phi}{2}}{\cos\theta_s}. \tag{31b}$$

The mirror surface (18a) is expressed in polar coordinates as

$$z_s = -\frac{r_s^2}{2R_m}\,\frac{\sin^2\theta_s\cos^2\frac{\phi}{2} + \cos^2\theta_s}{\cos\frac{\phi}{2}}. \tag{31c}$$

Applying similar computational procedure as in the previous subsection we obtain

$$R^{pq}_{mn} = C^{pq}_{mn}\, e^{i\psi^{pq}_{mn}} \cos\frac{\phi}{2}\int_0^{X_s} dX \left\{D^{pq}_{mn}(X)U^{pq}(X) + \xi \sin\frac{\phi}{2} E^{pq}(X)\left[V^{pq}_{mn}(X) + i\, W^{pq}_{mn}\right]\right\}, \tag{32a}$$

$$C_{mn}^{pq} = \frac{1}{2\pi} \left[\frac{m!n!\ \alpha^2}{(m+p)!(n+q)!}\right]^{1/2} , \tag{32b}$$

and the phase ψ_{mn}^{pq} is expressed by

$$\psi_{mn}^{pq} = (2m+p+1)\ \tan^{-1}\left(\frac{l_i}{b_i}\right) - (2n+q+1)\ \tan^{-1}\left(\frac{l_o}{b_o}\right) + k(l_i+l_o) \ . \tag{32c}$$

The integrals D^{pq}, E^{pq}, U_{mn}^{pq}, V_{mn}^{pq} and W_{mn}^{pq} are given by

$$D^{pq}(X) = \int_o^{2\pi} d\theta_m \ \frac{\cos p\left[\theta_i(\theta_s)\right]\ \cos q\left[\theta_o(\theta_s)\right]}{1 - \sin^2\frac{\phi}{2}\ \sin^2\theta_m} , \tag{33a}$$

$$E^{pq}(X) = \int_o^{2\pi} d\theta_m \ \frac{\sin\theta_m\ \cos p\left[\theta_i(\theta_s)\right]\ \cos q\left[\theta_o(\theta_s)\right]}{(1 - \sin^2\frac{\phi}{2}\ \sin^2\theta_m)^{3/2}} , \tag{33b}$$

$$U_{mn}^{pq}(X) = (\alpha^2 X)^{\frac{p}{2}}\ X^{\frac{q}{2}}\ L_m^p(\alpha^2 X)\ L_n^q(X)\ e^{-\frac{\alpha^2+1}{2}X} , \tag{33c}$$

$$V_{mn}^{pq}(X) = \frac{1}{\sqrt{2}} \left\{\frac{p-q}{2} - 3 - \frac{\alpha^2+1}{2}\ X + \left[\frac{L_m^{p\prime}(\alpha^2 X)}{L_m^p(\alpha^2 X)} - \frac{L_n^{q\prime}(X)}{L_n^q(X)}\right] X \right\}$$

$$(\alpha^2 X)^{\frac{p}{2}}\ X^{\frac{q+1}{2}}\ L_m^p(\alpha^2 X)\ L_n^q(X)\ e^{-\frac{\alpha^2+1}{2}X} , \tag{33d}$$

$$W_{mn}^{pq}(X) = \frac{1}{\sqrt{2}}\ (\alpha^2 X)^{\frac{p}{2}}\ X^{\frac{q+1}{2}}\ L_m^p(\alpha^2 X)\ L_n^q(X)\ e^{-\frac{\alpha^2+1}{2}X} . \tag{33e}$$

In obtaining (33a) - (33e), X was defined by $X = [1-\sin^2(\phi/2)\sin^2\theta]r^2/2W_o^{\ 2}$; thus, according to (22) and (30), the boundary X_s is $X_s = 2\rho^2/W_o^{\ 2}$. The lowest order terms for the first few elements are given in Appendix B for arbitrary deflection angle ϕ and $\alpha = 1$.

V. LIMITING CASES

When the mirror radius tends to infinity ($1/R_m \to 0$), or in cases of vertical incidence on the mirror ($\phi = 0$), the higher order corrections in the reflection matrix **R** disappear,

$$\mathbf{R} = \mathbf{R}(0) \tag{34}$$

in both representations. The nondiagonal elements in **R** stem from the finite mirror size only. If, in addition, the mirror size is very large, $\mu >> 1$, it is appropriate to take $W_o = W_i$ as best representation for the reflected radiation. The $\alpha = 1$ limit yields

$$R_{pq}^{mn} = \delta_{pq}^{mn}. \tag{35}$$

Thus, in case of large curved mirror and vertical incidence, or large plane mirror and arbitrary incidence, the reflection matrix is the identity matrix.

The case $\alpha = 1$ is of special interest for arbitrary angle of deflection ϕ and mirror curvature $1/R$, as it will be explained in the next section. For finite mirror size $\rho \geq W_o$, ($\mu \geq 1$), there exists zeroth order non-diagonal terms inside **R**(0). Since **R**(0) is independent of the angle of deflection ϕ, the finite mirror size yields the dominant contribution to the reflection into modes different than the incoming. The effects of the deflection of the light beam enter to first order in ξ, **R**(1), or higher. In the Hermite representation the elements $R_{pq}^{mn}(0)$ couple mode combinations with $m + p$ = even, $n + q$ = even. The elements with either $m + p$ or $n + q$ odd vanish because of the even/odd symmetry of the Hermite functions.

As the mirror size becomes very large and the limits of integration in (23) are extended, the orthogonality among the various modes $u_{mn}(r_s)$ becomes effective. The off-diagonal terms in **R**(0) become comparable to

the first order corrections roughly when $1/\mu^2 \sim \xi \sim \theta_d$. At the limit $\mu \to \infty$ all the nondiagonal elements of **R** are reduced to order ξ or higher,

$$R_{pq}^{mn} = \xi\, R_{pq}^{mn}(1) + 0(\xi^2), \qquad m \neq p,\ n \neq q, \tag{36a}$$

and the only matrix elements of zeroth order in ξ are the diagonal

$$R_{mn}^{mn} = R_{mn}^{mn}(0) + 0(\xi^2), \tag{36b}$$

in both Hermite and Laguerre represenations. The lowest correction in the diagonal elements is of second order ξ^2, while the first order contribution disappears. This is consistent with flux conservation during reflection in case of large mirror.

In obtaining Eqs. (28) and (32) it was assumed that $\Delta(x_s, y_s)$ is constant over S. According to (21) the variation of Δ is parametrized by $\xi^* = (kW_i{}^2/R_m)\ \xi$. When $(kW_i{}^2/R_m) \geq 1$, ξ^* becomes of order ξ and the effects of the slow phase slippage must be retained in (23). This effect, known as spherical aberration, causes additional corrections $\mathbf{R}^*(1)$, of order ξ^*,

$$\mathbf{R} = \mathbf{R}(0) + \xi\, \mathbf{R}(1) + \overset{*}{\xi}\ \overset{*}{\mathbf{R}}\ (\overset{*}{1}) + \ldots..$$

Spherical aberration does not disappear at the limit of large mirror size, as opposed to the effects discussed so far. In fact, when $\xi^* > \xi$, it places a lower limit on the off-diagonal terms in the reflection matrix,

$$R_{pq}^{mn} \geq \xi^*\ R_{pq}^{mn}(1)^*.$$

Perfect reflection, requiring $\xi^* = 0$, is possible only for plane mirror ($R_m \to \infty$) of large size.

The superposition principle can be used to describe reflection from more complex mirror surfaces. In case of a mirror with a hole the surface integral (14) over S_m is expressed as $\int_S = \int_{S1} - \int_{S2}$ where S_1 is defined by the mirror exterior boundary and S_2 is the surface of the hole. The total

reflection matrix **R** is given by $\mathbf{R} = \mathbf{R}(S_1) - \mathbf{R}(S_2)$, the difference in the reflection matrices associated with mirrors S_1 and S_2 respectively. The transmission matrix **T** through a screen with an aperture of area S is given by $\mathbf{T} = -\mathbf{R}$, **R** being the reflection matrix for a mirror matching the aperture S. The transmission matrix for radiation diffracted behind a finite size mirror is given by $\mathbf{T}' = \mathbf{1} - e^{i\pi}\,\mathbf{R}$ where **1** is the identity matrix.

VI. REFLECTION OF THE LOWEST ORDER MODE

The computation of all the truncated integrals for finite mirror surface is nontrivial. Most applications, however, involve the (0,0) lowest order mode as the dominant mode in both incoming and reflected radiation. The strategy here is to compute the element R_{00}^{00} of the reflection matrix first. Then the waist for the reflected modes W_o can be selected so that it maximizes R_{00}^{00}. The optimum representation condition

$$\frac{\partial R_{00}^{00}}{\partial \alpha} = 0, \tag{37}$$

puts the maximum amount of the reflected radiation in the lowest order mode (a different mode and matrix element may be chosen, if desired). It is pointed out that (37) does not improve the properties of the reflected radiation. It enables one to choose the best representation in terms of minimizing the coefficients of the undesired modes for the scattered radiation. Once W_o is fixed by (37) then the exact location and size of the waist(s) for the reflected modes is determined by solving the system of equations

$$\frac{1}{R_o} = \frac{l_o}{l_o^2 + b_o^2}, \tag{38a}$$

$$W_o = w_o \left[1 + \frac{l_o^2}{b_o^2}\right]^{1/2} . \tag{33b}$$

The element R_{00}^{00} is identical in both representations since the lowest order mode u_{00} is the same in rectangular and cylindrical coordinates. Performing the integration (29a) yields R_{00}^{00} to first order in ξ

$$R^{00}_{00} = \frac{2\alpha}{1+\alpha^2}\left[1 - e^{-(1+\alpha^2)\mu^2}\right] + O(\xi^2). \tag{39}$$

Note that the first order term vanishes and the lowest correction is of second order in ξ^2. The exact dependence on the mirror size ρ is parametrized by $\mu = \rho/W_o$, while $\alpha = W_i/W_o$ parametrizes the ratio of the incoming and scattered radiation spot sizes at the mirror. The optimization condition $\partial R^{00}_{00}(0)/\partial\alpha = 0$ yields, $\alpha^2 = 1 + \exp[-(1+\alpha^2)\mu^2][2\mu^2\alpha^4+(2\mu^2+1)\alpha^2-1]$. In case that the mirror cross section is much larger than the spot size of the incoming mode, $\mu >> 1$, $\alpha \to 1$ and the reflected spot size at the mirror matches the incoming, $W_o = W_i$.

Large mirror size is desired to maximize the total reflection coefficient. For incoming radiation of unity electromagnetic flux $P_i = |c^i|^2 = \Sigma \; |c^i{}_{pq}|^2 = 1$, the total reflection coefficient $\eta_R = P_o/P_i$ equals the reflected flux P_o,

$$P_o = |c^o|^2 = |\, R \cdot c^i |^2 = \sum_{mn} \sum_{pq} |R^{mn}_{pq} \; c^i{}_{pq}|^2. \tag{40}$$

In Fig. 3 we plot η_R for the lowest order incoming mode as a function of $\mu' = \cos(\phi/2)\; \rho/W_o = \cos(\phi/2)\; \mu$. μ' parametrizes the size of the mirror projection into the plane perpendicular to the incoming radiation direction. The incoming radiation has a wavelength $\lambda = 1\mu$ (10^{-4}cm), waist $w_i = 2\times10^{-1}$cm at distance $l_i = 1.8\times10^2$cm from the mirror and radius of curvature (at the mirror) $R_i = 8.95\times10^3$cm. The mirror has a radius of curvature $R_m = 8.95\times10^3$cm, yielding reflected modes of $R_o = 8.95\times10^3$ (again l_o and w_o depend on the choice of W_o). In Fig. 4 we plot the magnitude of the reflection coefficients $|R^{00}_{pq}|$ of the lowest order mode (0,0) into the firstr 25 modes (p,q) with $p \leq q \leq 5$, as a function of μ'. The deflection angle is 90^o and the ratio of the spot

sizes is 1. Increasing mirror size maximizes the diagonal element and minimizes scattering into other modes. The spherical aberration was retained inside (23) in evaluating the matrix elements. Its effect is small, since for the above parameters $\xi^* = 0.28\xi$, and a good agreement is observed with the constant Δ theoretical limit. In particular, the dominant off-diagonal terms couple the (0,0) incoming mode to the (1,0), (3,0) and (3,2) reflected modes only, according to the selection rules, Eqs. A(10). Comparing Figs. 3 and 4 with the next plots shows that the relative mirror size to the radiation spot size is the most important parameter to determine the reflection into other than the incoming modes.

In Fig. 5 we fix the mirror size $\mu' = 2$ and the angle $\phi = 90^o$ and vary the spot size ratio α. The best representation, maximizing R_{00}^{00} and minimizing R_{pq}^{00} is obtained at $\alpha = 1$. However, for small mirror $\mu' = 0.66$, the maximum for R_{00}^{00} occurs at $\alpha \simeq 0.70$ (see Fig. 6). Radiation reflected off mirrors smaller than the incoming spot size is best described by outgoing modes of reduced spot size $W_o < W_i$. Also note from Fig. 6b that for small mirror size the total power reflected into the first 25 modes never exceeds 80% of the incoming flux; even with many more modes η_R remains less than 1. In Fig. 7 the reflection coefficients R_{pq}^{00} are plotted as functions of the angle of deflection ϕ for fixed $\alpha = 1$, $\mu' = 2$. It is seen that, for sufficiently large reflecting surface and good choice of the spot size W_o, the reflection matrix is not very sensitive to ϕ and the off-diagonal terms remain small.

The main conclusions so far are summarized as follows. When the mirror size is ≥ 2.5 times the incoming spot size, the fraction of the incident power scatterd into different modes is of order ξ^2 for

$kW^2_i/R_m < 1$, or $(\xi^*)^2$ for $kW^2_i/R_m > 1$. This holds for a wide range of deflection angles ϕ. It will be shown in the next section that cross-polarization effects are of the same order. In most applications both ξ and ξ^* are less than 10^{-2}. To this end, scattering losses will be smaller than the losses caused by the finite reflectivity (i.e., absorption) by the mirror, for most dielectrics.

VII. CROSS-POLARIZATION EFFECTS

The curvature of the mirror surface produces a cross-coupling between the transverse components of the incoming and reflected radiation. Inserting expressions (21) for the normal unit vector to the mirror inside the boundary conditions Eq. (7a), the full source term $\mathbf{A}^s = (A_x{}^s, A_y{}^s, A_z{}^s)$ for an incoming wave $\mathbf{A}^i = (A_x{}^i, A_y{}^i, 0)$ is given by

$$A_x{}^s = - A_x{}^i + 2 \tan\frac{\phi}{2} \frac{y_s}{R_m} A_y{}^i ,$$

$$A_y{}^s = - A_y{}^i + 2 \tan\frac{\phi}{2} \frac{y_s}{R_m} A_x{}^i ,$$

$$A_z{}^s = 2 \cos\frac{\phi}{2} \frac{x_s}{R_m} A_x{}^i + 2 \frac{y_s}{R_m} . \tag{41}$$

In the above relations, the components of $\mathbf{A}^i$ and $\mathbf{A}^s$ are given in coordinate systems aligned with the incoming and outgoing radiation, respectively. According to (41) the reflection of a plane polarized wave generates components polarized in every direction, including A_z. These cross polarization effects enter to order ξ and result in a small rotation of the polarization angle.

The radiation steming from the $A_z{}^s$ component will propagate perpendicularly to the direction of interest $\hat{z}_o$ and escapes the resonator as pure reflection loss. The relation between the incoming and reflected transverse components, including cross-polarization effects, is now given by

$$\begin{pmatrix} C_x{}^o \\ C_y{}^o \end{pmatrix} = \begin{pmatrix} R & Q \\ Q & R \end{pmatrix} \begin{pmatrix} C_x{}^i \\ C_y{}^i \end{pmatrix} . \tag{42}$$

The matrix $\mathcal{R}$ has been computed in the previous section. Substitution

of the additional cross-terms in Eq. (41) inside the propagator integral (6) yields

$$Q_{(xy)pq}^{mn} = Q_{(yx)pq}^{mn} = Q_{pq}^{mn} ,$$

where

$$Q_{pq}^{mn} = \iint_S dx_s dy_s \; 2 \tan\frac{\phi}{2} \frac{y_s}{R_m} \frac{\bar{u}_{mn}(x_s,y_s)\bar{u}_{pq}(x_s,y_s)}{\left[1 + \frac{l_o{}^2}{b_o{}^2}\right]^{1/2}} \left[\frac{1 + \frac{z_o{}^2(x_s,y_s)}{b_o{}^2}}{1 + \frac{z_i{}^2(x_s,y_s)}{b_i{}^2}}\right]^{1/2}$$

$$\times \quad e^{i\delta_{pq}^{i}(x_s,y_s) - i\delta_{mn}^{o}(x_s,y_s)} \left[\cos\frac{\phi}{2}\left(1 - \frac{x_s}{R_m}\sin\frac{\phi}{2} - \frac{x_s{}^2\sin^2\frac{\phi}{2}}{R_m{}^2}\right)\right] . \tag{43}$$

In Gaussian-Hermite representation, we obtain

$$Q_{pq}^{mn}(1) = \sqrt{2}\xi \, C_{pq}^{mn} \, e^{i\psi_{pq}^{mn}} \tan\frac{\phi}{2} \, G_{pq}^{mn} , \tag{44a}$$

with

$$G_{pq}^{mn} = \int_{-X_s}^{X_s} dX \int_{-Y_s}^{Y_s} dY \; H_p(\alpha X)\; H_q(Y)\; Y\; H_m(X)\; H_n(Y)\; e^{-\frac{\alpha^2+1}{2}(X^2+Y^2)} . \tag{44b}$$

In Gaussian-Laguerre representation, we have

$$Q_{mn}^{pq}(1) = \sqrt{2}\;\xi\;\sin\frac{\phi}{2} \int_o^{X_s} dX \; G_{mn}^{pq}(X)\; B^{pq}(X), \tag{45a}$$

where

$$G_{mn}^{pq}(X) = (\alpha^2 X)^{\frac{1}{2}} \, X^{\frac{q+1}{2}} \, L_m^p(\alpha^2 X)\; L_n^q(X)\; e^{-\frac{\alpha^2+1}{2}X} , \tag{45b}$$

and

$$B^{pq}(X) = \int_0^{2\pi} d\theta \; \frac{\cos\theta \cos\left[p\theta_i(\theta)\right]\cos\left[q\theta_i(\theta)\right]}{1 - \sin^2\frac{\phi}{2}\sin^2\theta} . \tag{45c}$$

In both representations, cross polarization effects enter to order ξ. In case of vertical incidence ($\phi = 0$) with arbitrary curvature $1/R_m$, or plane mirror ($\xi \sim 1/R_m = 0$) and arbitrary incidence ϕ, Q goes to zero. Transverse vector components are reflected independently of each other in these two limits. Some of the elements of Q (in both representations) are given in Appendix C for large ($\rho \gg W_i$) mirror.

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Appendix A. Computation of the Hermitian Matrix Elements.

The integrals (29) will be evaluated here in case the mirror size ρ is much larger than the incoming mode spot size W_i, $\rho \cos \phi/2 >> W_i$. Then the limits of the surface integrals can be extended to infinity, and the spot size for the outgoing modes W_o matches that of the incoming at the mirror, i.e., $\alpha = 1$. We use the notation

$$\psi_n = e^{-\frac{X^2}{2}} H_n(X), \tag{A1}$$

the recurrence relation

$$H_n'(X) = 2n\, H_{n-1}(X), \tag{A2}$$

and the orthonormality properties

$$\int_{-\infty}^{\infty} dX\ \psi_n(X)\psi_m(X) = 2^n n! \sqrt{\pi}\ \delta_{m,n}, \tag{A3}$$

$$\int_{-\infty}^{\infty} dX\ \psi_n(X)\ X\ \psi_m(X) \equiv \langle X\rangle_{m,n} = \sqrt{\pi}\left(2^{n-1} n!\,\delta_{m,n-1} + 2^n (n+1)!\,\delta_{m,n+1}\right), \tag{A4}$$

$$\int_{-\infty}^{\infty} dX\ \psi_n(X)\ X^2\ \psi_m(X) \equiv \langle X^2\rangle_{m,n}$$

$$= \sqrt{\pi}\left\{2^{n-2} n!\,\delta_{m,n-2} + 2^{n-1}(2n+1)n!\,\delta_{m,n} + 2^n (n+2)!\,\delta_{m,n+2}\right\}, \tag{A5}$$

to obtain

$$I_{pq}^{mn} = \pi\ 2^{m+n} n! m!\ \delta_{p,m}\delta_{q,n}, \tag{A6}$$

$$M_{pq}^{mn} = -\frac{3}{2}\langle X\rangle_{p,m}\left(\sqrt{\pi}\ 2^n n!\right)\delta_{q,n},$$

$$- \frac{\sqrt{2}}{4} \left[2k\langle X^2\rangle_{p-1,m} - 2m \langle X^2\rangle_{m-1,p}\right]\left(\sqrt{\pi}\ 2^n n!\right)\delta_{q,n}$$

$$- \frac{\sqrt{2}}{4} \left[2k\ 2^m m!\ \delta_{p-1,m} - 2m\ 2^{m-1}(m-1)!\,\delta_{p,m-1}\right]\langle Y^2\rangle_{q,n}, \quad \text{(A7)}$$

$$N_{pq}^{mn} = \frac{1}{\sqrt{2}} \left[(m+n+1)\frac{b_i R_m}{1_i R_i} + (p+q+1)\frac{b_o R_m}{1_o R_o}\right] \langle X\rangle_{p,m}\left(\sqrt{\pi}2^n n!\,\xi\right)\delta_{q,n}. \quad \text{(A8)}$$

Inserting expressions (A3)-(A5) into Eqs. (A6)-(A8) we obtain

$$R_{pq}^{mn}(0) = \delta_{p,m}\ \delta_{q,n}, \quad \text{(A9)}$$

$$R_{pq}^{mn}(1) = \quad \text{(A10)}$$

$$\tan\frac{\phi}{2} \left\{\left\{-\left[\frac{3}{2}\left(\frac{m-1}{2}\right)^{1/2} + \frac{1}{\sqrt{2}}\left(\frac{m-1}{2}\right)^{3/2} - \left(\frac{2m-1}{4} + \frac{2n+1}{4}\right)m^{1/2}\right]\delta_{p,m-1}\right.\right.$$

$$- \left[\frac{3}{2}\left(\frac{m+1}{2}\right)^{1/2} - \frac{1}{\sqrt{2}}\left(\frac{m+1}{2}\right)^{3/2} + \frac{1}{\sqrt{2}}\left(\frac{m+1}{2}\right)^{1/2}\left(\frac{2m+1}{2} + \frac{2n+1}{2}\right)\right]\delta_{p,m+1}$$

$$\left. - \frac{1}{\sqrt{2}}\left(\frac{(m+1)(m+2)(m+3)}{8}\right)^{1/2}\delta_{p,m+3} + \frac{1}{\sqrt{2}}\left(\frac{m(m-1)(m-2)}{8}\right)^{1/2}\delta_{p,m-3}\right\}\delta_{q,n}$$

$$- \frac{1}{\sqrt{2}}\left[\left(\frac{m+1}{2}\right)^{1/2}\delta_{p,m+1} - \left(\frac{m}{2}\right)^{1/2}\delta_{p,m-1}\right]\left(\frac{(n-1)(n-2)}{4}\right)^{1/2}\delta_{q,n-2}$$

$$- \frac{1}{\sqrt{2}}\left[\left(\frac{m+1}{2}\right)^{1/2}\delta_{p,m+1} - \left(\frac{m}{2}\right)^{1/2}\delta_{p,m-1}\right]\left(\frac{(n+1)(n+2)}{4}\right)^{1/2}\delta_{q,n+2}$$

$$+$$

$$\left.\frac{i}{\sqrt{2}}\left[(m+n+1)\frac{b_i R_m}{1_i R_i} + (p+q+1)\frac{b_o R_m}{1_o R_o}\right]\left[\left(\frac{m-1}{2}\right)^{1/2}\delta_{p,m-1} + \left(\frac{m+1}{2}\right)^{1/2}\delta_{p,m+1}\right]\delta_{q,n}\right\}.$$

$\mathbf{R}$ is diagonal to zeroth order. The lowest order correction couples m with $m\pm1$, $m\pm3$ in the X-direction and n with n, $n\pm2$ in the Y-direction. The reflection matrix is not symmetric, $R^{mn}_{pq} \neq R^{pq}_{mn}$. Also, it is not invariant to interchanging X and Y. This means that the modes $u_{mn}(x,y)$ and $u_{nm}(x,y)$ with $m \neq n$ are reflected differently.

Appendix B. Computation of the First Matrix Elements in Laguerre Representation.

Representation using Gaussian-Laguerre modes may be advantageous in numerical simulations because fewer Laguerre modes than Hermite modes are required to represent close-to-axisymmetric radiation profiles with the same accuracy. However, the computation of Eqs. (33a) to (33e) is not so straightforward. The integrations (33a) and (33b) for I^{pq} and K^{pq} over the polar angle θ_s involve trigonometric functions of complicated arguments $\theta_i(\theta_s)$ and $\theta_o(\theta_s)$, given implicitly by Eq. (31b). There is no simple recurrence formula for this calculation. The first few elements are computed here by expansions in powers of $r_s/R_m < \zeta \sim \xi$. Substituting from (31b) inside (33) and renormalizing $r_s^{\,2}/2W_o^{\,2} = X/(1-\sin^2\phi/2\sin^2\theta)$, one obtains, to first order in ξ,

$$D^{00}(X) = \int_o^{2\pi} d\theta \, \frac{1}{1-\sin^2\frac{\phi}{2}\sin^2\theta}\,, \tag{B0}$$

$$D^{10}(X) = -I^{01}(X) = \frac{1}{2\sqrt{2}}\,\xi\, X \tan\frac{\phi}{2}\int_o^{2\pi} d\theta \, \frac{\cos^2\theta}{\left(1-\sin^2\frac{\phi}{2}\sin^2\theta\right)^2}, \tag{B1}$$

$$D^{11}(X) = \int_o^{2\pi} d\theta \, \frac{\cos^2\theta}{\left(1-\sin^2\frac{\phi}{2}\sin^2\theta\right)^2}, \tag{B2}$$

$$D^{-1-1}(X) = \cos^2\frac{\phi}{2}\int_o^{2\pi} d\theta \, \frac{\sin^2\theta}{\left(1-\sin^2\frac{\phi}{2}\sin^2\theta\right)^2}, \tag{B3}$$

$$D^{-10}(X) = -D^{0-1}(X) = -\frac{1}{2\sqrt{2}}\,\xi\, X\tan\frac{\phi}{2}\int_o^{2\pi} d\theta \, \frac{\left(1-\cos^2\frac{\phi}{2}\tan^2\theta\right)\cos^2\theta}{\left(1-\sin^2\frac{\phi}{2}\sin^2\theta\right)^2}, \tag{B4}$$

$$D^{-11}(X) = D^{1-1}(X) = 0. \qquad \text{(B5)}$$

We only need E^{pq} to zeroth order in ξ, obtaining

$$E^{-10}(X) = -E^{0-1}(X) = \int_0^{2\pi} d\theta \, \frac{\cos\frac{\phi}{2}\sin^2\theta}{\left(1-\sin^2\frac{\phi}{2}\sin^2\theta\right)^{3/2}}, \qquad \text{(B6)}$$

and

$$E^{pq} = 0 + O(\xi) \qquad \text{for } (p,q) \neq (-1,0), (0-1). \qquad \text{(B7)}$$

The integrals (B1) - (B7) are evaluated using the formula

$$\int_0^{\pi/2} dx \, \frac{\sin^{\mu}x \cos^{\nu}x}{\left(1-k^2\sin^2 x\right)^{\rho}} = \frac{1}{2} B\left(\frac{\mu+1}{2}, \frac{\nu+1}{2}\right) F\left(\rho, \frac{\mu+1}{2}, \frac{\mu+\nu+2}{2}, k^2\right), \qquad \text{(B8)}$$

where $B(p,q) = \Gamma(p)\Gamma(q)/\Gamma(p+q)$, Γ is the factorial function and F is the hypergeometric function. The radial integrations for U, V and W are performed directly, using the expressions $L_m^p(x)$ for the Laguerre functions and the identities

$$\int_0^{\infty} e^{-x} x^{-1/2} dx = \sqrt{\pi},$$

$$\int_0^{\infty} e^{-x} x^n dx = n!,$$

$$\int_0^{\infty} e^{-x} x^{n+1/2} dx = \frac{1}{2} \cdot \frac{3}{2} \ldots (n+1/2)\sqrt{\pi}. \qquad \text{(B9)}$$

Again, we extend the limits of integration to infinity assuming $\rho \cos(\phi/2) >> W_i$ and $\alpha = 1$. The zeroth order contribution is given by

$$R_{00}^{00}(0) = 1,$$
$$R_{11}^{00}(0) = 1,$$

$$R^{-1-1}_{\ 1\ 1}(0) = \cos^3 \frac{\phi}{2}\, F\left(2,\frac{3}{2},2,\sin^2 \frac{\phi}{2}\right),$$

$$R^{11}_{11}(0) = \cos \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\sin^2 \frac{\phi}{2}\right),$$

$$R^{11}_{00}(0) = \cos \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\sin^2 \frac{\phi}{2}\right),$$

$$R^{-1-1}_{\ 0\ 0}(0) = \cos^2 \frac{\phi}{2}\, F\left(2,\frac{3}{2},2,\sin^2 \frac{\phi}{2}\right). \tag{B10}$$

The first order corrections in ξ are given by

$$R^{10}_{00}(1) = -R^{01}_{00}(1) = \left(\frac{\pi}{2}\right)^{1/2} \frac{3}{8} \sin \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right),$$

$$R^{10}_{01}(1) = -R^{01}_{10}(1) = \left(\frac{\pi}{2}\right)^{1/2} \frac{9}{16} \sin \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right),$$

$$R^{00}_{01}(1) = R^{00}_{10}(1) = 0,$$

$$R^{10}_{10}(1) = -R^{01}_{01}(1) = \left(\frac{\pi}{2}\right) \quad \frac{3^{1/2}}{16} \sin \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right),$$

$$R^{10}_{11}(1) = -R^{01}_{11}(1) = -\left(\frac{\pi}{2}\right)^{1/2} \frac{39}{32} \sin \frac{\phi}{2}\, F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right), \tag{B11}$$

and

$$R^{-10}_{\ 00}(1) = \sin \frac{\phi}{2} \left\{ \pm \left[\left(\frac{\pi}{2}\right)^{1/2} \frac{3}{8} F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right) - \cos^2 \frac{\phi}{2}\, F\left(2,\frac{3}{2},2,\ \sin^2 \frac{\phi}{2}\right) \right] \right.$$

$$\left. \pm \frac{1}{2} \cos^2 \frac{\phi}{2}\, F\left(\frac{3}{2},\frac{3}{2},2,\ \sin^2 \frac{\phi}{2}\right) \begin{bmatrix} -4.5 & + i2^{-1/2} \\ -3.5 & + i2^{-1/2} \end{bmatrix} \right\},$$

$$R^{-10}_{\ 01}(1) = \sin \frac{\phi}{2} \left\{ \pm \left[\left(\frac{\pi}{2}\right)^{1/2} \frac{9}{16} F\left(2,\frac{1}{2},2,\ \sin^2 \frac{\phi}{2}\right) - \cos^2 \frac{\phi}{2}\, F\left(2,\frac{3}{2},2,\ \sin^2 \frac{\phi}{2}\right) \right] \right.$$

$$\left. \pm \frac{1}{2} \cos^2 \frac{\phi}{2}\, F\left(\frac{3}{2},\frac{3}{2},2,\ \sin^2 \frac{\phi}{2}\right) \left[2 - i\ 2^{-1/2}\right] \right\},$$

$$R^{-10}_{11}(1) = \sin\frac{\phi}{2}\left\{\pm\left[\left(\frac{\pi}{2}\right)^{1/2}\frac{39}{32}F\left(2,\frac{1}{2},2,\ \sin^2\frac{\phi}{2}\right) - \cos^2\frac{\phi}{2}F\left(2,\frac{3}{2},2,\ \sin^2\frac{\phi}{2}\right)\right],\right.$$

$$\left.\pm\frac{1}{2}\cos^2\frac{\phi}{2}F\left(\frac{3}{2},\frac{3}{2},2,\ \sin^2\frac{\phi}{2}\right)\begin{bmatrix}4.5 + i2^{-1/2}\\ 3.5 + i2^{-1/2}\end{bmatrix}\right\}. \tag{B12}$$

The (-) sign and the lowest row inside the last square bracket in (B12) correspond to exchanging indices,

$$R^{pq}_{mn} \leftrightarrow R^{qp}_{nm}.$$

Note also that the elements $R^{1-1}_{m\ n}$, R^{-11}_{mn}, coupling sine and cosine modes, are of order ξ^2 or higher for every m, n,

$$R^{-1\ 1}_{\ m\ n} \sim O(\xi^2). \tag{B13}$$

Appendix C. Computation of the Cross-Polarization Matrix Elements

We compute here some of the first order cross-polarization matrix elements in case of large mirror size $\rho \gg W_i$ and $\alpha \to 1$. In the Gaussian-Hermite representation we find from (44b), using the notation of Appendix A,

$$G_{pq}^{mn} = 2^m m! \sqrt{\pi} \langle Y \rangle_{q,n} \delta_{p,m}$$

$$= \pi\, 2^m m! \left\{ 2^{n-1} n!\, \delta_{q,n-1} + 2^n (n+1)!\, \delta_{q,n+1} \right\} \delta_{p,m}, \tag{C1}$$

yielding

$$Q_{pq}^{mn}(1) = \xi \sin\frac{\phi}{2} \left\{ \sqrt{n}\, \delta_{q,n-1} + \sqrt{n+1}\, \delta_{q,n+1} \right\} \delta_{p,m}. \tag{C2}$$

In Gaussian-Laguerre representation we only have to compute $B^{pq}(X)$, Eq. (45c), to zeroth order in ξ. Applying the methods of Appendix B, we find

$$B^{pq} = O(\xi), \text{ if } p,q \neq (1,0), (0,1),$$

$$B^{10} = B^{01} = \int_0^{2\pi} \frac{d\theta \cos^2\theta}{\left(1 - \sin^2\frac{\phi}{2}\sin^2\theta\right)^{3/2}}. \tag{C3}$$

Noting that $G_{mn}(X)$ is the same as $W_{mn}^{pq}(X)$, Eq. (33e), and inserting (C3) and (33e) inside (45a), we obtain

$$Q_{10}^{10}(1) = Q_{01}^{01}(1) = 0,$$

$$Q_{00}^{10}(1) = Q_{00}^{01}(1) = \xi \sin\frac{\phi}{2}\, F\left(\frac{3}{2}, \frac{1}{2}, 2, \sin^2\frac{\phi}{2}\right),$$

$$Q_{01}^{10}(1) = Q_{10}^{01}(1) = -\xi \sin\frac{\phi}{2}\, F\left(\frac{3}{2}, \frac{1}{2}, 2, \sin^2\frac{\phi}{2}\right),$$

$$Q_{11}^{10}(1) = Q_{11}^{01}(1) = \sqrt{2}\, \xi \sin\frac{\phi}{2}\, F\left(\frac{3}{2}, \frac{1}{2}, 2, \sin^2\frac{\phi}{2}\right). \tag{C4}$$

Appendix D. Small Aperture Limit

We have seen in Sec. V that in case of mirror surface S_1 with an aperture of area S_2 the reflection matrix is given by

$$R(S;W_o) = R(S_1;W_o) - R(S_2;W_o) \tag{D1}$$

In case that $S_2 << S_1$ the spot size W_o optimizing the representation for the scattered radiation will be determined predominantly by the surface S_1. Thus, the formula (23) with W_o given from

$$\frac{\partial R(S_1;\alpha)}{\partial\alpha} = 0, \tag{D2}$$

can be used for the modal decomposition of the scattered radiation. According to Eqs. (28a) and (29a) for the Hermite representation, and Eqs. (32a) and (33c) for the Laguerre representation, the lowest order contribution from a small aperture $\rho_2 << W_o$ scales as $R(S_2;W_o) \sim \xi^2$.

In some cases, however, it is important to know the <u>total</u> radiation diffracted through a small hole, rather than the modal decomposition. In case of small apperture ρ_2,

$$\rho_2{}^2 << l_o k^{-1} \quad \text{or} \quad \lambda >> \rho_2^2/l_o, \tag{D3}$$

where $l_o \sim z$ is the distance of the observation point from the mirror, the paraxial approximation, Eq. (8b) is taken one step further, setting

$$\frac{k}{z-z_o}\left[\left(x-x_o\right)^2 + \left(y-y_o\right)^2\right] \simeq \frac{k}{z-z_o}\left[\left(x^2+y^2\right) - 2xx_o - 2yy_o\right]. \tag{D4}$$

Substituting (D4) inside (8a) we obtain the "far field" limit of the diffracted radiation

$$A^o(x,y,z=0) = \frac{ik}{2\pi z_o}\int dx_o \int dy_o A^i(x_o,y_o)(\hat{n}\cdot\delta\hat{r})$$

$$\exp\left\{ -\frac{ik}{2z_o}\left[\left(x^2+y^2\right) - 2xx_o - 2yy_o\right]\right\}, \tag{D5}$$

also known as Fraunhoffer diffraction. The condition (D3) can only be valid for appertures much smaller than the spot size W_o at the mirror, $\rho_2 << W_o$, since for $\rho_2 \sim W_o$ (D3) is violated, $k\rho_2^{\ 2} \sim kW_o^{\ 2} \sim kw_o^{\ 2}\ (1+l_o^{\ 2}/b_o^{\ 2}) \sim b_o(1+l_o^{\ 2}/b_o^{\ 2}) > l_o$. Neglecting terms of order $kx_m^{\ 2}/l_o \sim k\rho_2^{\ 2}/l_o << 1$ means that terms of order x_s/l_m, $x_s/l_o << 1/kx_s$, where $kx_s > 1$, must also be neglected. The source term can be written as $A^i[x_i,y_i] \simeq A^i[x_s,y_s]$. Rescaling variables to

$$K_x = \cos\frac{\phi}{2}\frac{kx}{l_o}, \qquad K_y = \frac{ky}{l_o} \tag{D6}$$

we obtain

$$A^o(x,y,0) = e^{i\left(pq_o + k\frac{x^2+y^2}{l_o}\right)} \cos\frac{\phi}{2}\frac{ik}{2\pi l_o}\int dx_s \int dy_s A^i(x_s,y_s)e^{iK_xx_s+iK_yy_s}. \tag{D7}$$

According to (D7) the outgoing radiation is the Fourier transform of the incoming radiation in respect to K_x, K_y. Defining the "polar" coordinates $K = (K_x^{\ 2}+K_y^{\ 2})^{1/2}$, $\Theta = \tan^{-1}(K_x/K_y)$, we obtain, for $A^i(x_s,y_s) = \sum_{k,l} C^i_{pq}\ u_{pq}\ (x_s,y_s)$, the scattered radiation

$$A^o(x,y,0) = \sum_{p,q} c^i_{pq}\ \frac{ika_{pq}}{2\pi l_o}\ e^{ik\left(l_o + \frac{x^2+y^2}{2l_o}\right)}\ \text{X}$$

$$\cos\frac{\phi}{2}\int dx_s\int dy_s\; H_p\left[\sqrt{2}\,\frac{x_s}{W_i}\cos\frac{\phi}{2}\right]H_q\left[\sqrt{2}\,\frac{y_s}{W_i}\right]\, e^{-\frac{x_s^{\,2}\cos^2\frac{\phi}{2}+y_s^{\,2}}{W_i^{\,2}}}$$

$$\sum_{m=-\infty}^{\infty}\;\sum_{n=-\infty}^{\infty} J_n(Kx_s)\; J_m(Ky_s)\; e^{i(m-n)\theta}\; i^{m+n}\; e^{in\frac{\pi}{2}}. \qquad \text{(D8)}$$

The zeroth order contribution in $x_s/W_i \ll 1$ yields

$$A^o(x,y,0) \simeq c^i_{oo}\,\frac{ika_{oo}}{2\pi l_o}\; e^{ik\left(l_o+\frac{x^2+y^2}{2l_o}\right)}$$

$$X\,\cos\frac{\phi}{2}\int dx_s\int dy_s\; J_o(Kx_s)\; J_o(Ky_s)e^{-\frac{x_s^{\,2}\cos^2\frac{\phi}{2}+y_s^{\,2}}{W_i^{\,2}}}\quad . \qquad \text{(D9)}$$

The waist size w_f for the Fraunhoffer modes is given by the zeros of the Bessel functions

$$K(w_f)\; x_s \sim \frac{kw_f}{l_o}\,\rho_2 \sim 2\pi. \qquad \text{(D10)}$$

Therefore, the diffraction angle θ_f is

$$\theta_f \simeq \frac{w_f}{l_o} \sim \frac{\lambda}{\rho_2}. \qquad \text{(D11)}$$

The requirement $\theta_f \ll 1$ for the validity of the paraxial approximation puts a lower limit in the aperture size ρ_2

$$\rho_2 \gg \lambda. \qquad \text{(D12)}$$

In case the aperture size is of the order of the wavelength λ the scattered wavefunctions are spherical rather than Gaussian. Because the overall effect of a scatterer with size $\rho_2 \sim \lambda$ is very small, the familiar from quantum mechanics Born approximation, involving perturbation theory, is applicable in that case.

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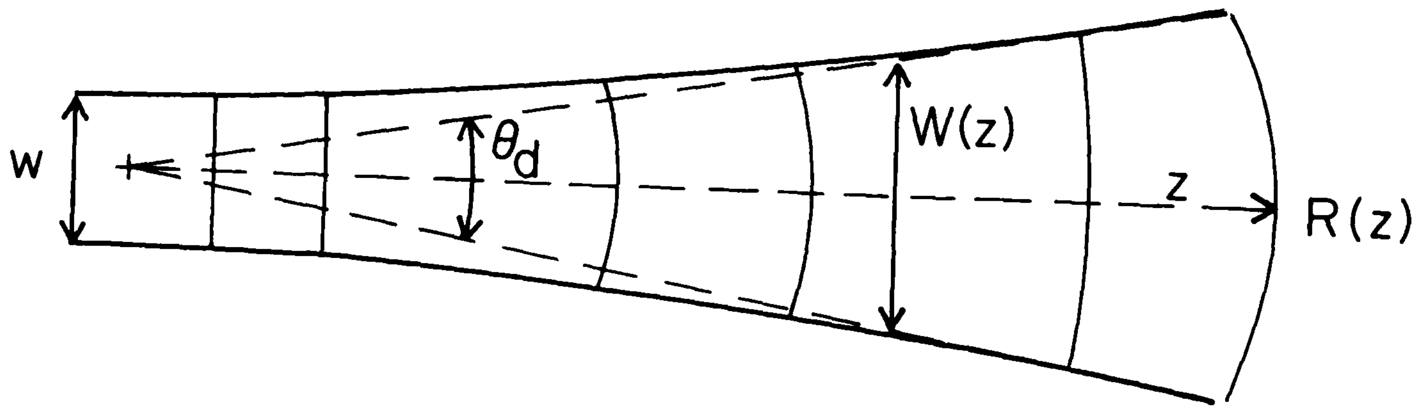


Figure 1 Schematic illustration of the radiation envelope for a Gaussian eigenmode.

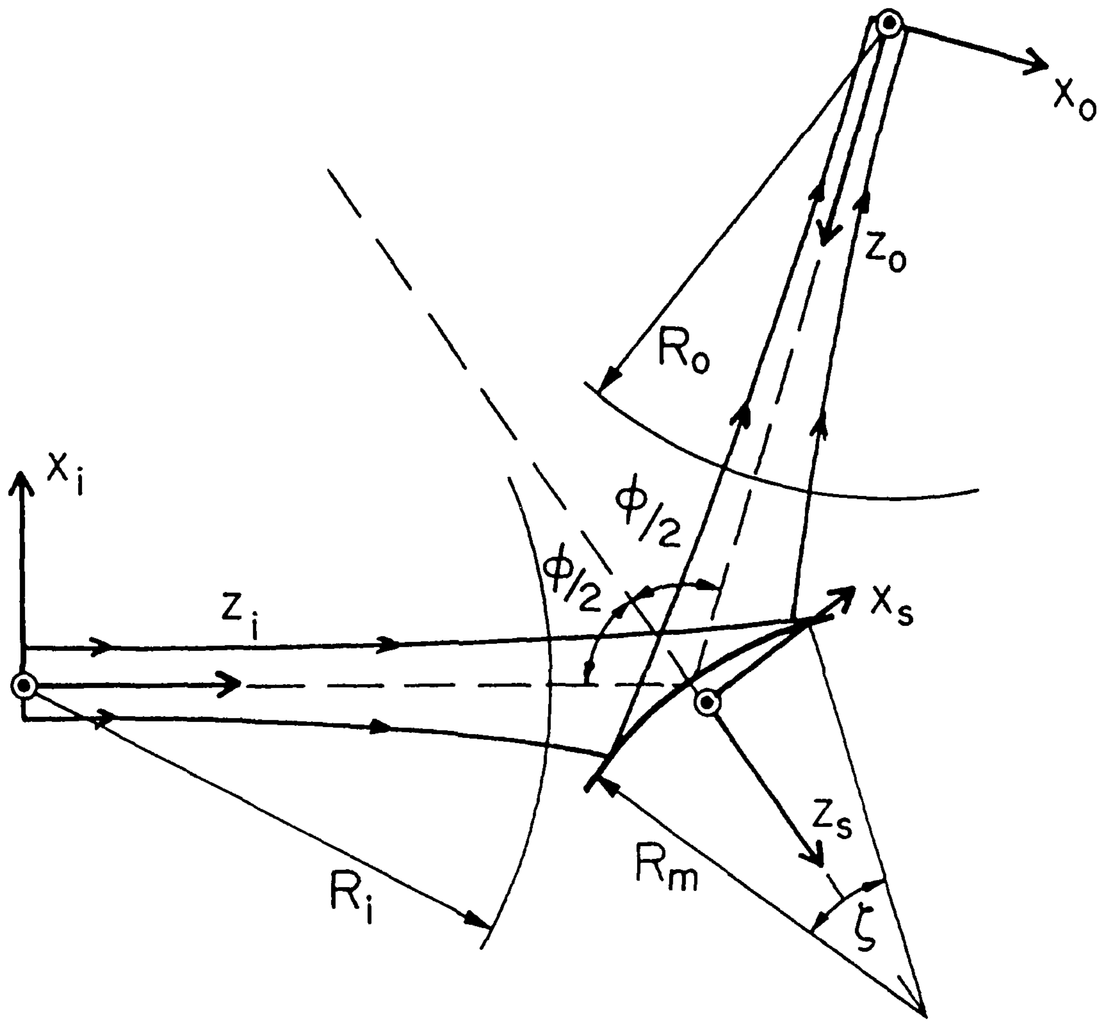


Figure 2 Reflection geometry.

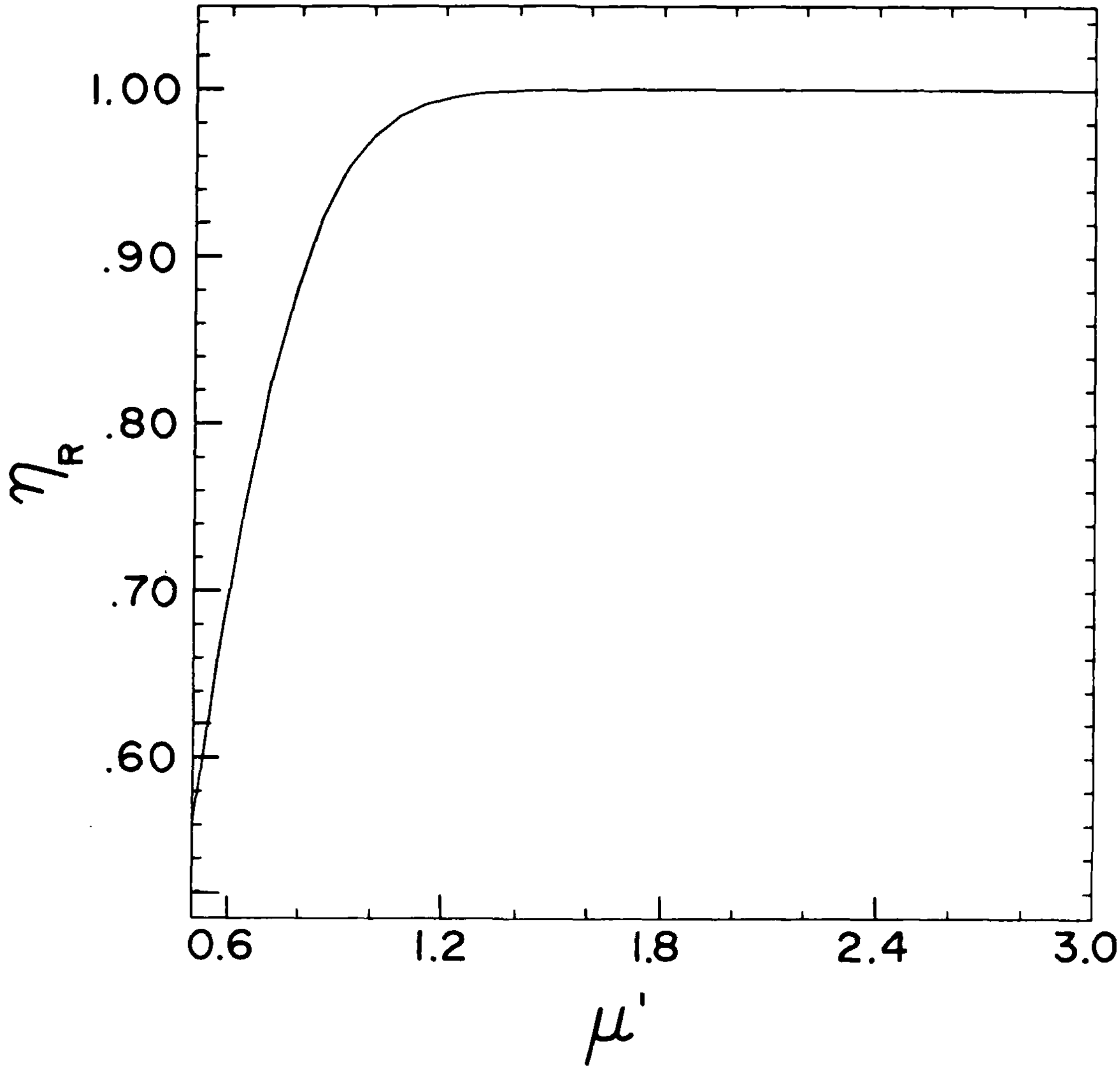


Figure 3 Plot of the total reflection coefficient η_R for the lowest order mode as a function of the mirror size μ' for $\phi = 90°$. The radiation has wavelength $\lambda = 10^{-4}$cm, waist $w_i = 2\times10^{-1}$cm at distance $l_i = 1.8\times10^{2}$cm from the mirror and radius of curvature $R_i = R_m = 8.95\times10^{3}$cm.

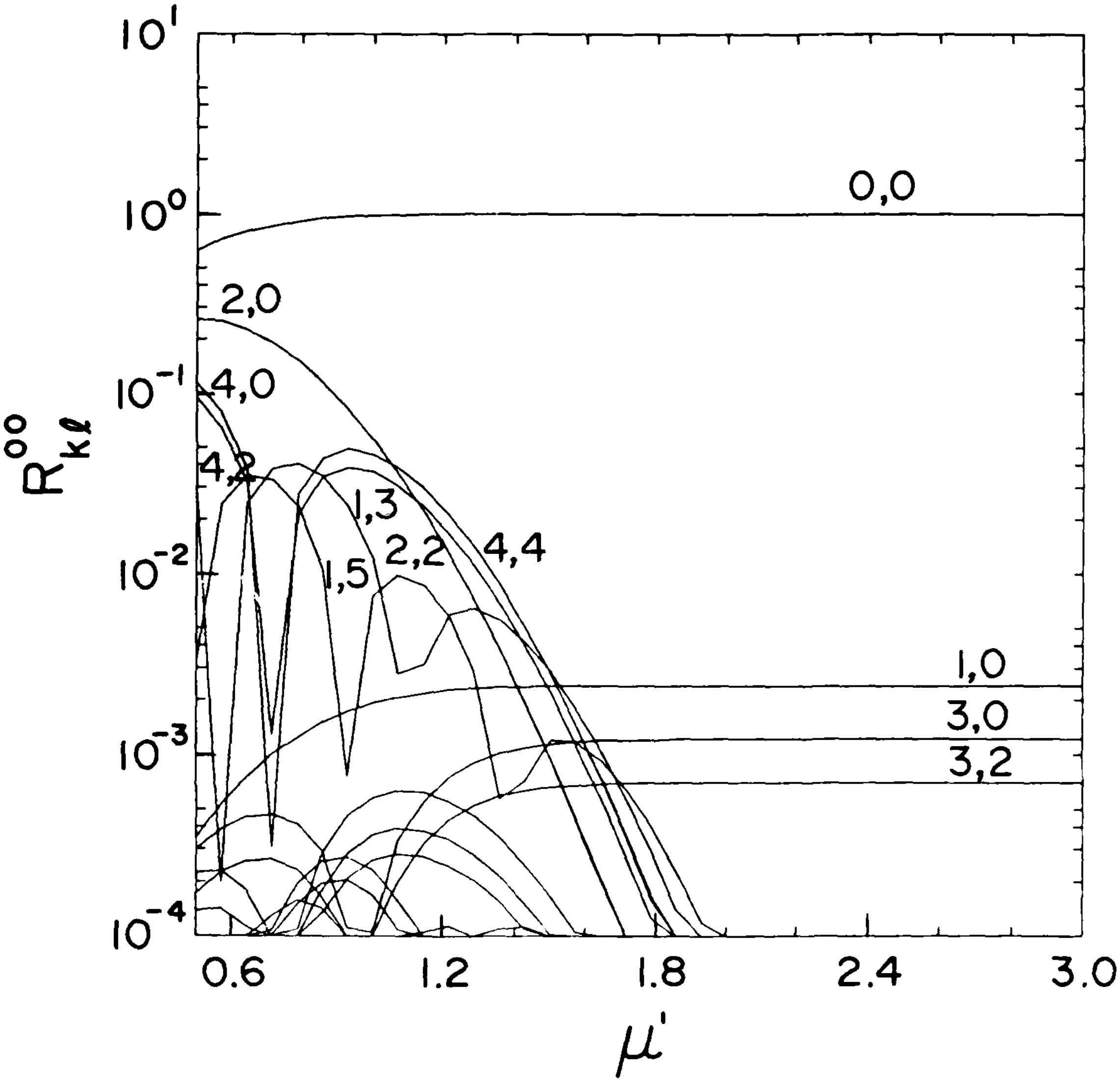
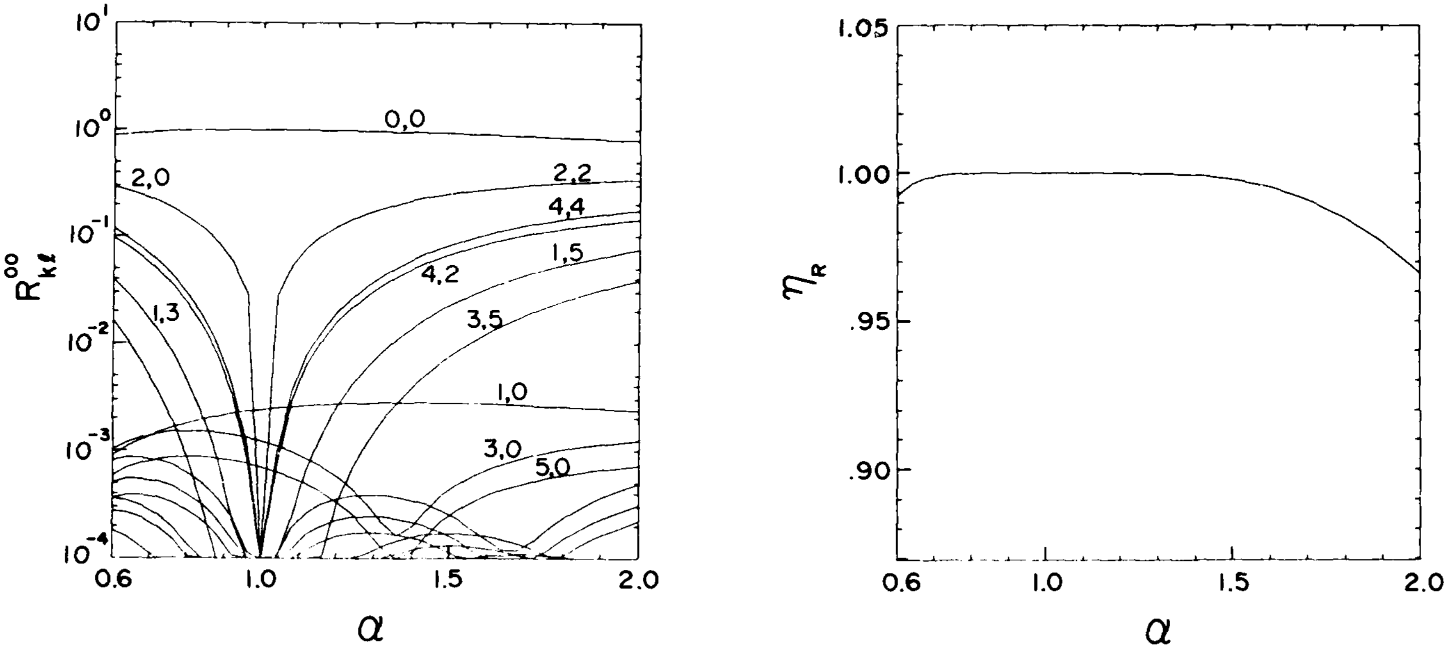


Figure 4 Reflection matrix elements for the lowest order mode (0,0) into the first 25 modes (p,q) against the relative mirror size μ'. The magnitude $|R^{00}_{pq}|$ is plotted for angle of deflection $\phi = 90°$, $\alpha = 1$ ($W_i = W_o$).



(a) (b)

Figure 5 Plots of the reflection matrix elements $|R^{00}_{pq}|$ against the spot size ratio α for $\mu' = 2$ and $\phi = 90°$. Radiation parameters are the same as in Fig. 3.

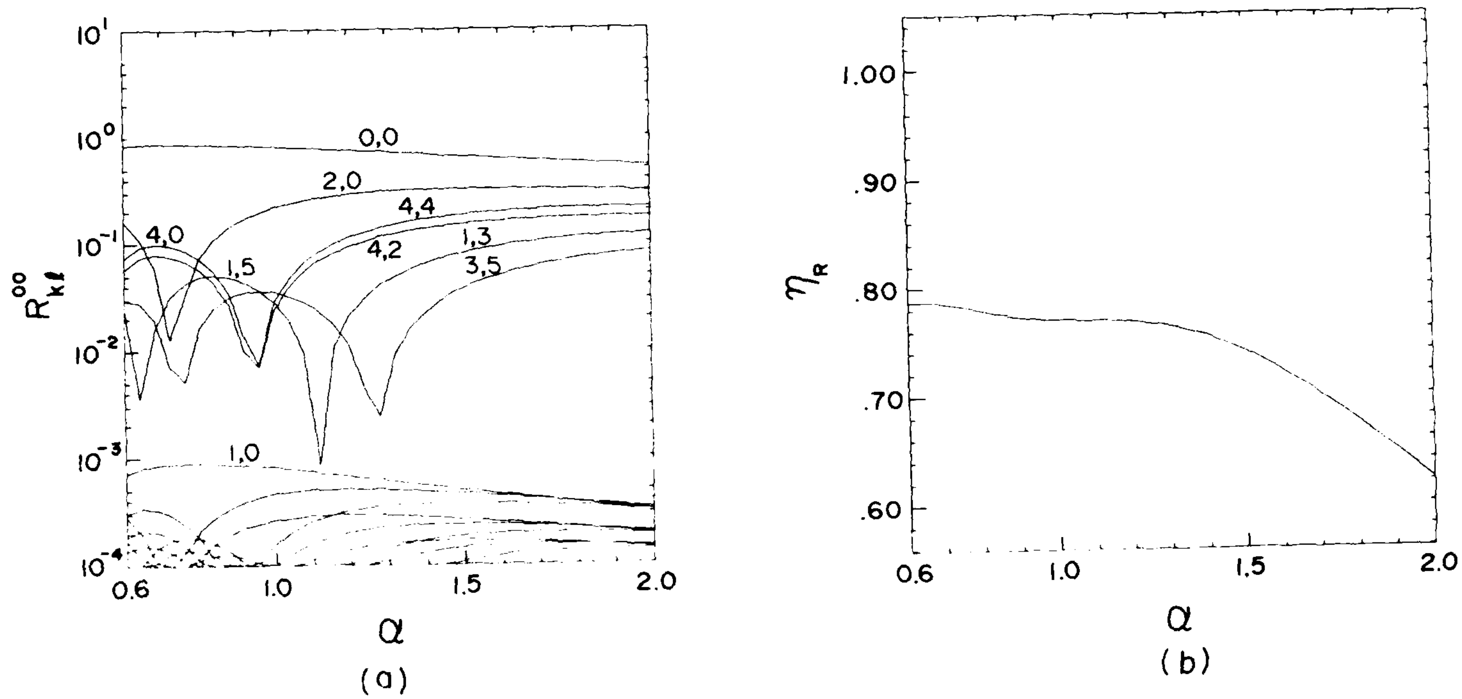


Figure 6 Same as in Fig. 5 for μ' = 0.66.

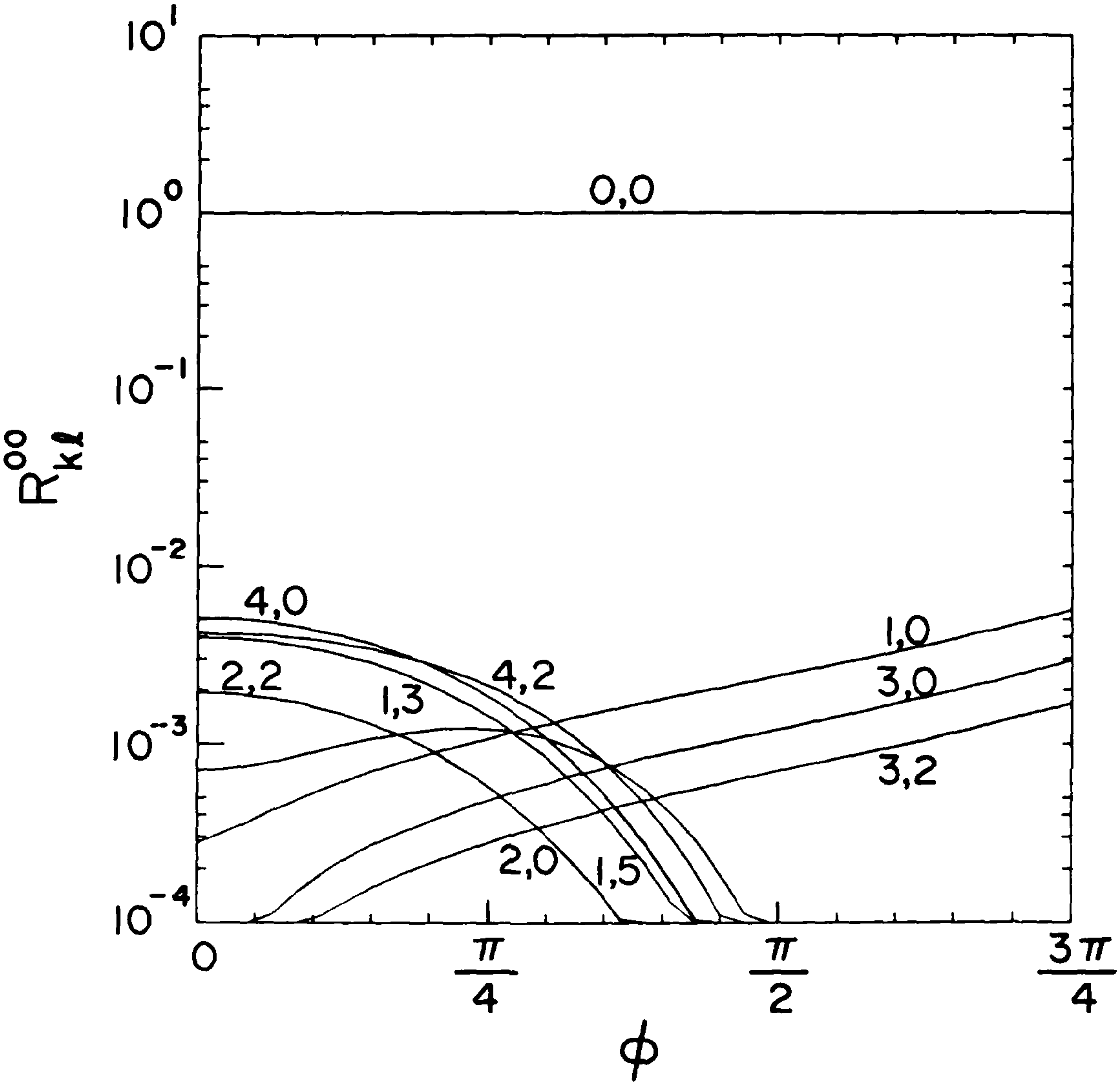


Figure 7 Plots of the reflection matrix elements $|R^{00}_{pq}|$ against the angle of deflection ϕ for $\mu' = 2$ and $\alpha = 1$.

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